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Generalized Universality for TMD Distribution Functions

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Azimuthal asymmetries in high-energy processes, most pronounced showing up in combination with single or double (transverse) spin asymmetries, can be understood with the help of transverse momentum dependent (TMD) parton distribution and fragmentation functions. These appear in correlators containing expectation values of quark and gluon operators. TMDs allow access to new operators as compared to collinear (transverse momentum integrated) correlators. These operators include nontrivial process dependent Wilson lines breaking universality for TMDs. Making an angular decomposition in the azimuthal angle, we define a set of universal TMDs of definite rank, which appear with process dependent gluonic pole factors in a way similar to the sign of T-odd parton distribution functions in deep inelastic scattering or the Drell-Yan process. In particular, we show that for a spin 1/2 quark target there are three pretzelosity functions.

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1. Introduction

To study the connection between partons and hadrons in high energy processes through parton distribution functions (PDF) and parton fragmentation functions (PFF), the starting points are forward matrix elements of parton fields, such as the quark-quark correlator

$$\Phi_{ij}(p|p) = \int \frac{d^4\xi}{(2\pi)^4} e^{i p \cdot \xi} \langle P | \bar{\psi}_j(0) \psi_i(\xi) | P \rangle, \quad (1)$$

where a summation over color indices is understood. This replaces the correlator $\Phi \propto (\not{p} + m)$ for a single incoming fermion. In the case of hadrons one also needs quark-quark-gluon correlators such as

$$\Phi_{Aij}^\mu(p - p_1, p_1|p) = \int \frac{d^4\xi d^4\eta}{(2\pi)^8} e^{i(p-p_1) \cdot \xi} e^{i p_1 \cdot \eta} \langle P | \bar{\psi}_j(0) A^\mu(\eta) \psi_i(\xi) | P \rangle. \quad (2)$$

2 *M. G. A. Buffing and P. J. Mulders*

The basic idea is to factorize these hadronic (soft) parts in a full diagrammatic approach and parametrize them in terms of PDFs. This requires high energies in which case the momenta of different hadrons obey $P \cdot P' \propto Q^2$, where Q^2 is the hard scale in the process. In that case the hadronic momenta can be treated as light-like vectors P and the hard process brings in a conjugate light-like vector n such that $P \cdot n = 1$, for instance $n = P'/P \cdot P'$. One makes a Sudakov expansion of the parton momenta,

$$p = xP + p_T + (p \cdot P - xM^2)n, \quad (3)$$

with $x = p^+ = p \cdot n$. In any contraction with vectors outside the correlator, the component xP contributes at order Q , the transverse component at order $M \sim Q^0$ and the remaining component contributes at order M^2/Q . This allows consecutive integration of the components to obtain from the fully unintegrated correlator in Eq. (1) the TMD light-front (LF) correlator

$$\Phi_{ij}(x, p_T; n) = \int \frac{d\xi \cdot P d^2\xi_T}{(2\pi)^3} e^{i p \cdot \xi} \langle P | \bar{\psi}_j(0) \psi_i(\xi) | P \rangle \Big|_{\xi \cdot n = 0}, \quad (4)$$

the collinear light-cone (LC) correlator

$$\Phi_{ij}(x) = \int \frac{d\xi \cdot P}{2\pi} e^{i p \cdot \xi} \langle P | \bar{\psi}_j(0) \psi_i(\xi) | P \rangle \Big|_{\xi \cdot n = \xi_T = 0 \text{ or } \xi^2 = 0}, \quad (5)$$

or the local matrix element

$$\Phi_{ij} = \langle P | \bar{\psi}_j(0) \psi_i(\xi) | P \rangle \Big|_{\xi=0}. \quad (6)$$

The importance of integrating at least the light-cone (minus) component $p^- = p \cdot P$ is that the expression is at equal time, i.e. time-ordering is not relevant anymore for TMD or collinear PDFs¹. For local matrix elements one can calculate the anomalous dimensions, which show up as the Mellin moments of the splitting functions that govern the scaling behavior of the collinear correlator $\Phi(x)$. We note that the collinear correlator is not simply an integrated TMD. The dependence on an upper limit $\Phi(x; Q^2) = \int^Q d^2p_T \Phi(x, p_T)$ is found from the anomalous dimensions (splitting functions). One has an α_s/p_T^2 behavior of TMDs that is calculable using collinear TMDs and which matches to the intrinsic non-perturbative p_T -behavior². We note that in operator product expansion language, the collinear correlators involve operators of definite twist, while TMD correlators involve operators of various twist.

2. Color gauge invariance

In order to determine the importance of a particular correlator in a hard process, one can do a dimensional analysis to find out when they contribute in an expansion in the inverse hard scale. Dominant are the ones with lowest canonical dimension obtained by maximizing contractions with n , for instance for quark or gluon fields the

minimal canonical dimensions $\dim[\bar{\psi}(0)\not{n}\psi(\xi)] = \dim[F^{n\alpha}(0)F^{n\beta}(\xi)] = 2$, while an example for a multi-parton combination gives $\dim[\bar{\psi}(0)\not{n}A_T^\alpha(\eta)\psi(\xi)] = 3$. Equivalently, one can maximize the number of P 's in the parametrization of Φ_{ij} . Of course one immediately sees that any number of collinear $n \cdot A(\eta) = A^n(\eta)$ fields doesn't matter. Furthermore one must take care of color gauge invariance, for instance when dealing with the gluon fields and one must include derivatives in color gauge invariant combinations. With dimension zero there is $iD^n = i\partial^n + gA^n$ and with dimension one there is $iD_T^\alpha = i\partial_T^\alpha + gA_T^\alpha$. The color gauge-invariant expressions for quark and gluon distribution functions actually include gauge link operators,

$$U_{[0,\xi]} = \mathcal{P} \exp \left(-i \int_0^\xi d\zeta_\mu A^\mu(\zeta) \right), \quad (7)$$

connecting the nonlocal fields,

$$\Phi_{qij}^{[U]}(x, p_T; n) = \int \frac{d\xi \cdot P d^2\xi_T}{(2\pi)^3} e^{ip \cdot \xi} \langle P | \bar{\psi}_j(0) U_{[0,\xi]} \psi_i(\xi) | P \rangle \Big|_{LF}, \quad (8)$$

$$\begin{aligned} \Gamma_g^{[U,U']}{}^{\mu\nu}(x, p_T; n) &= \int \frac{d\xi \cdot P d^2\xi_T}{(2\pi)^3} e^{ip \cdot \xi} \\ &\times \text{Tr} \langle P, S | F^{n\mu}(0) U_{[0,\xi]} F^{n\nu}(\xi) U'_{[\xi,0]} | P, S \rangle \Big|_{LF}. \end{aligned} \quad (9)$$

For transverse separations, the gauge links involve paths running along the minus direction to $\pm\infty$ (dimensionally preferred), which are closed with one or more transverse pieces at light-cone infinity. The two simplest possibilities are $U^{[\pm]} = U_{[0,\pm\infty]}^n U_{[0_T,\xi_T]}^T U_{[\pm\infty,\xi]}^n$, leading to gauge link dependent quark TMDs $\Phi_q^{[\pm]}(x, p_T)$ ^{3,4}. For gluons, the correlator involves color gauge-invariant traces of field-operators $F^{n\alpha}$, which are written in the color-triplet representation, requiring the inclusion of *two* gauge links $U_{[0,\xi]}$ and $U'_{[\xi,0]}$. Again the simplest possibilities are the past- and future-pointing gauge links $U^{[\pm]}$, giving even in the simplest case four gluon TMDs $\Gamma^{[\pm,\pm^\dagger]}(x, p_T)$.

Using the dimensional analysis to collect the leading contributions in an expansion in the inverse hard scale, one will need the above quark and gluon TMDs for the description of azimuthal dependence. Taking the Drell-Yan (DY) process as an example, one can look at the cross section depending on the (small!) transverse momentum q_T of the produced lepton pair,

$$\begin{aligned} \sigma(x_1, x_2, q_T) &= \int d^2p_{1T} d^2p_{2T} \delta^2(p_{1T} + p_{2T} - q_T) \\ &\times \Phi_1^{[-]}(x_1, p_{1T}) \bar{\Phi}_2^{[-^\dagger]}(x_2, p_{2T}) \hat{\sigma}(x_1, x_2, Q), \end{aligned} \quad (10)$$

which involves a convolution of TMDs. What is more important, it is the color flow in the process, in this case neutralized in the initial state, that determines the path

4 *M. G. A. Buffing and P. J. Mulders*

in the gauge link in the TMDs, in this case past-pointing ones. In contrast in semi-inclusive deep inelastic scattering one finds that the relevant TMD is $\Phi^{[+]}$ with a future-pointing gauge link. In a general process one can find more complex gauge links including besides Wilson line elements also Wilson loops. In particular when the transverse momentum of more than one hadron is involved, such as e.g. in the DY case above, it may be impossible to have just a single TMD for a given hadron because color gets entangled^{5,6}.

The correlators including a gauge link can be parametrized in terms of TMD PDFs^{7,8} depending on x and p_T^2 ,

$$\begin{aligned} \Phi^{[U]}(x, p_T; n) = & \left\{ f_1^{[U]}(x, p_T^2) - f_{1T}^{\perp[U]}(x, p_T^2) \frac{\epsilon_T^{p_T S_T}}{M} + g_{1s}^{[U]}(x, p_T) \gamma_5 \right. \\ & \left. + h_{1T}^{[U]}(x, p_T^2) \gamma_5 \not{S}_T + h_{1s}^{\perp[U]}(x, p_T) \frac{\gamma_5 \not{p}_T}{M} + i h_1^{\perp[U]}(x, p_T^2) \frac{\not{p}_T}{M} \right\} \frac{\not{P}}{2}, \end{aligned} \quad (11)$$

with the spin vector parametrized as $S^\mu = S_L P^\mu + S_T^\mu + M^2 S_L n^\mu$ and shorthand notations for $g_{1s}^{[U]}$ and $h_{1s}^{\perp[U]}$,

$$g_{1s}^{[U]}(x, p_T) = S_L g_{1L}^{[U]}(x, p_T^2) - \frac{p_T \cdot S_T}{M} g_{1T}^{[U]}(x, p_T^2). \quad (12)$$

For quarks, these include not only the functions that survive upon p_T -integration, $f_1^q(x) = q(x)$, $g_1^q(x) = \Delta q(x)$ and $h_1^q(x) = \delta q(x)$, which are the well-known collinear spin-spin densities (involving quark and nucleon spin) but also momentum-spin densities such as the Sivers function $f_{1T}^{\perp q}(x, p_T^2)$ (unpolarized quarks in a transversely polarized nucleon) and spin-spin-momentum densities such as $g_{1T}(x, p_T^2)$ (longitudinally polarized quarks in a transversely polarized nucleon).

The parametrization for gluons, following the naming convention in Ref. 9, is given by

$$\begin{aligned} 2x \Gamma^{\mu\nu[U]}(x, p_T) = & -g_T^{\mu\nu} f_1^{g[U]}(x, p_T^2) + g_T^{\mu\nu} \frac{\epsilon_T^{p_T S_T}}{M} f_{1T}^{\perp g[U]}(x, p_T^2) \\ & + i \epsilon_T^{\mu\nu} g_{1s}^{g[U]}(x, p_T) + \left(\frac{p_T^\mu p_T^\nu}{M^2} - g_T^{\mu\nu} \frac{p_T^2}{2M^2} \right) h_1^{\perp g[U]}(x, p_T^2) \\ & - \frac{\epsilon_T^{p_T \{\mu} p_T^{\nu\}}}{2M^2} h_{1s}^{\perp g[U]}(x, p_T) - \frac{\epsilon_T^{p_T \{\mu} S_T^{\nu\}} + \epsilon_T^{S_T \{\mu} p_T^{\nu\}}}{4M} h_{1T}^{g[U]}(x, p_T^2). \end{aligned} \quad (13)$$

3. Transverse moments

In many cases, it is convenient to construct moments of TMDs in the same way as one considers moments of collinear functions. For $\Phi(x)$ in Eq. (5) one constructs moments

$$\begin{aligned} x^N \Phi(x) = & \int \frac{d\xi \cdot P}{2\pi} e^{i p \cdot \xi} \langle P | \bar{\psi}(0) (i \partial^n)^N U_{[0, \xi]}^n \psi(\xi) | P \rangle \Big|_{LC} \\ = & \int \frac{d\xi \cdot P}{2\pi} e^{i p \cdot \xi} \langle P | \bar{\psi}(0) U_{[0, \xi]}^n (i D^n)^N \psi(\xi) | P \rangle \Big|_{LC}. \end{aligned} \quad (14)$$

Integrating over x one finds the connection of the Mellin moments of PDFs with local matrix elements with specific anomalous dimensions, which via an inverse Mellin transform define the splitting functions. Similarly one can consider transverse moment weighting starting with the light-front TMD in Eq. (4),

$$p_T^\alpha \Phi^{[\pm]}(x, p_T; n) = \int \frac{d\xi \cdot P d^2\xi_T}{(2\pi)^3} e^{i p \cdot \xi} \times \langle P | \bar{\psi}(0) U_{[0, \pm\infty]}^n U_{[0_T, \xi_T]}^T i D_T^\alpha(\pm\infty) U_{[\pm\infty, \xi]}^n \psi(\xi) | P \rangle \Big|_{LF}. \quad (15)$$

Integrating over p_T gives the lowest transverse moment, which appears in the q_T -weighted result of Eq. (10). This moment involves twist-3 (or higher) collinear multi-parton correlators, in particular the quark-quark-gluon correlator

$$\Phi_F^{n\alpha}(x - x_1, x_1 | x) = \int \frac{d\xi \cdot P d\eta \cdot P}{(2\pi)^2} e^{i(p-p_1) \cdot \xi} \times e^{i p_1 \cdot \eta} \langle P | \bar{\psi}(0) U_{[0, \eta]}^n F^{n\alpha}(\eta) U_{[\eta, \xi]}^n \psi(\xi) | P \rangle \Big|_{LC}. \quad (16)$$

In terms of this correlator and the similarly defined correlator $\Phi_D^\alpha(x - x_1, x_1 | x)$ one finds

$$\int d^2 p_T p_T^\alpha \Phi^{[U]}(x, p_T) = \tilde{\Phi}_\partial^\alpha(x) + C_G^{[U]} \pi \Phi_G^\alpha(x), \quad (17)$$

with

$$\begin{aligned} \tilde{\Phi}_\partial^\alpha(x) &= \Phi_D^\alpha(x) - \Phi_A^\alpha(x) \\ &= \int dx_1 \Phi_D^\alpha(x - x_1, x_1 | x) - \int dx_1 \text{PV} \frac{1}{x_1} \Phi_F^{n\alpha}(x - x_1, x_1 | x), \\ \Phi_G^\alpha(x) &= \Phi_F^{n\alpha}(x, 0 | x). \end{aligned}$$

The latter is referred to as a gluonic pole or ETQS-matrix element^{10,11}. They are multiplied with gluonic pole factors $C_G^{[U]}$ (e.g. $C_G^{[\pm]} = \pm 1$), that tell us that new functions are involved with characteristic process dependent behavior^{12,13}. This behavior is for the single transverse moments also coupled to the behavior under time-reversal. While $\tilde{\Phi}_\partial^\alpha$ is T-even, Φ_G^α is T-odd. Since time-reversal is a good symmetry of QCD, the appearance of T-even or T-odd functions in the parametrization of the correlators is linked to specific observables with this same character. In particular single spin asymmetries are T-odd observables.

The weighting with transverse momenta can also be analyzed by studying the parametrization in PDFs. For single p_T -weighting, only PDFs with one prefactor of p_T in the parametrization in Eq. (11) survive. The $\tilde{\Phi}_\partial^\alpha(x)$ matrix element receives contributions from T-even PDFs, while the $\Phi_G^\alpha(x)$ matrix element receives contributions from T-odd PDFs, see Ref. 14 for a detailed study of this. For the single weighted results, thus, the behavior under time-reversal can be used to identify the

6 *M. G. A. Buffing and P. J. Mulders*

process dependent parts and we find

$$f_{1T}^{\perp(1)[U]}(x) = C_G^{[U]} f_{1T}^{\perp(1)}(x), \quad (18)$$

where transverse weighting for PDFs is defined as

$$f_{\dots}^{(n)}(x) = \int d^2 p_T \left(\frac{-p_T^2}{2M^2} \right)^n f_{\dots}(x, p_T^2) \quad (19)$$

for weighting with n transverse moments. The fact that depending on the process $\Phi^{[\pm]}$ are the correlators to be used, leads to the sign change¹⁵ for the T-odd functions in such processes or to more complex factors if more complex gauge links are involved¹⁶. The importance of Eq. (18) is the appearance of a universal function with calculable process (link) dependent numbers rather than many process dependent functions that are somehow related. For gluon TMDs, there are already for single weighting two functions $\Gamma_G^{(f/d)}$ and hence two different gluonic pole factors $C_G^{[U](f/d)}$, because there are two ways to construct color singlets from the (in that case) three gluon fields that are involved using the f_{abc} or d_{abc} structure constants. The appearance of two different gluon Sivers functions was pointed out in Ref. 17.

The situation with universality for fragmentation functions¹⁸ is easier because the gluonic pole matrix elements vanish in that case^{19,20,21}. Nevertheless, there exist T-odd fragmentation functions, but their QCD operator structure is T-even. These T-odd functions then appear in the parametrization of $\tilde{\Phi}_{\partial}^{\alpha}$. Hence, there is no process dependence from gluonic pole factors.

The use of transverse moments in the description of azimuthal asymmetries via transverse momentum weighting of the cross section can be extended to higher moments involving higher harmonics such as $\cos(2\varphi)$. Also here process dependence may come in from double gluonic pole matrix elements $\Phi_{GG}^{\alpha\beta}$, which are twist four operators. This affects studies that involve the quark TMD $h_{1T}^{\perp q}(x, p_T)$ (pretzelosity distribution) in Eq. (11) or the gluon Boer-Mulders function $h_1^{\perp g}(x, p_T)$ (linearly polarized gluons in unpolarized targets) in Eq. (13). For instance, for quarks one finds for the simplest gauge links,

$$h_{1T}^{\perp(2)[\pm]}(x) = h_{1T}^{\perp(2)(A)}(x) + h_{1T}^{\perp(2)(B1)}(x), \quad (20)$$

where the functions $h_{1T}^{\perp(2)(A)}(x)$, $h_{1T}^{\perp(2)(B1)}(x)$ are universal. Actually the latter function corresponds to a correlator Φ_{GG} , involving a color structure $\text{Tr}_c [FF\psi\bar{\psi}]$. For more complex gauge links one actually needs a third (universal) pretzelosity second transverse moment because there is another possible color structure²².

4. TMDs of definite rank

An interesting possibility to obtain universal TMDs is to start with a parametrization that involves the symmetric traceless tensors $p_T^{\alpha_1 \dots \alpha_m}$ of rank m , such as

$$p_T^{\alpha}, \quad p_T^{\alpha\beta} = p_T^{\alpha} p_T^{\beta} - \frac{1}{2} p_T^2 g_T^{\alpha\beta} \dots \quad (21)$$

Depending on the rank different correlators come in, involving operator combinations of gluons, covariant derivatives and A -fields. Minimizing the twist we have

$$\begin{aligned}\Phi^{[U]}(x, p_T) &= \Phi(x, p_T^2) + \pi C_G^{[U]} \frac{p_{Ti}}{M} \Phi_G^i(x, p_T^2) + \pi^2 C_{GG,c}^{[U]} \frac{p_{Tij}}{M^2} \Phi_{GG,c}^{ij}(x, p_T^2) + \dots \\ &+ \frac{p_{Ti}}{M} \tilde{\Phi}_\partial^i(x, p_T^2) + \pi C_G^{[U]} \frac{p_{Tij}}{M^2} \tilde{\Phi}_{\{\partial G\}}^{ij}(x, p_T^2) + \dots \\ &+ \frac{p_{Tij}}{M^2} \tilde{\Phi}_{\partial\partial}^{ij}(x, p_T^2) + \dots, \end{aligned} \quad (22)$$

with a summation over the color structures c . The reproduction of the transverse moments provides the proper identification of universal TMD functions,

$$\Phi(x, p_T^2) = \left\{ f_1(x, p_T^2) + S_L g_1(x, p_T^2) \gamma_5 + h_1(x, p_T^2) \gamma_5 \not{S}_T \right\} \frac{\not{P}}{2}, \quad (23)$$

$$\frac{p_{Ti}}{M} \tilde{\Phi}_\partial^i(x, p_T^2) = \left\{ h_{1L}^\perp(x, p_T^2) S_L \frac{\gamma_5 \not{P}_T}{M} - g_{1T}(x, p_T^2) \frac{p_T \cdot S_T}{M} \gamma_5 \right\} \frac{\not{P}}{2}, \quad (24)$$

$$\frac{p_{Ti}}{M} \Phi_G^i(x, p_T^2) = \frac{1}{\pi} \left\{ -f_{1T}^\perp(x, p_T^2) \frac{\epsilon_T^{\rho\sigma} p_{T\rho} S_{T\sigma}}{M} + i h_{1T}^\perp(x, p_T^2) \frac{\not{P}_T}{M} \right\} \frac{\not{P}}{2}, \quad (25)$$

$$\frac{p_{Tij}}{M^2} \tilde{\Phi}_{\partial\partial}^{ij}(x, p_T^2) = h_{1T}^{\perp(A)}(x, p_T^2) \frac{p_{Tij} S_T^i \gamma_5 \gamma_T^j}{M^2} \frac{\not{P}}{2}, \quad (26)$$

$$\frac{p_{Tij}}{M^2} \Phi_{GG,1}^{ij}(x, p_T^2) = \frac{1}{\pi^2} h_{1T}^{\perp(B1)}(x, p_T^2) \frac{p_{Tij} S_T^i \gamma_5 \gamma_T^j}{M^2} \frac{\not{P}}{2}, \quad (27)$$

$$\frac{p_{Tij}}{M^2} \Phi_{GG,2}^{ij}(x, p_T^2) = \frac{1}{\pi^2} h_{1T}^{\perp(B2)}(x, p_T^2) \frac{p_{Tij} S_T^i \gamma_5 \gamma_T^j}{M^2} \frac{\not{P}}{2}, \quad (28)$$

$$\frac{p_{Tij}}{M^2} \tilde{\Phi}_{\{\partial G\}}^{ij}(x, p_T^2) = 0. \quad (29)$$

We note that the rank zero functions in Eq. (23) depend on x and p_T^2 and involve traces, to be precise $g_1(x, p_T^2) = g_{1L}^{[U]}(x, p_T^2)$ and $h_1(x, p_T^2) = h_{1T}^{[U]}(x, p_T^2) - (p_T^2/2M^2) h_{1T}^{\perp[U]}(x, p_T^2)$. As remarked before, for the pretzelosity there are three universal functions with in general

$$h_{1T}^{\perp[U]}(x, p_T^2) = h_{1T}^{\perp(A)}(x, p_T^2) + C_{GG,1}^{[U]} h_{1T}^{\perp(B1)}(x, p_T^2) + C_{GG,2}^{[U]} h_{1T}^{\perp(B2)}(x, p_T^2). \quad (30)$$

For the simplest gauge links we have $C_{GG,1}^{[\pm]} = 1$ and $C_{GG,2}^{[\pm]} = 0$, which shows e.g. that $h_{1T}^{\perp[\text{SIDIS}]}(x, p_T^2) = h_{1T}^{\perp[\text{DY}]}(x, p_T^2)$, but that for other processes (with more complicated gauge links) other combinations of the three possible pretzelosity functions occur. For a spin 1/2 target the above set of TMDs is complete. There are no higher rank functions. For a spin 1 target^{23,24} and for gluons, there are higher rank functions^{22,25}. For the fragmentation correlator there is for rank 2 only a single (T-even) pretzelosity function $H_{1T}^\perp(z, k_T^2)$ appearing in the parametrization of the correlator $\Delta_{\partial\partial}^{\alpha\beta}(x, p_T^2)$. The rank of the various correlators is shown in Table 1 with the results for nucleon TMD PDFs summarized in Tables 2 (unpolarized) and 3 (polarized) and those for nucleon TMD PFFs in Tables 4 and 5.

Table 1. The correlators of definite rank in the full TMD correlator, ordered in columns according to the number of gluonic poles (G) and ordered in rows according to the number of contributing partial derivatives ($\partial = D - A$). The rank of these operators is equal to the sum of these numbers. Their twist is equal to the rank + 2.

GLUONIC POLE RANK			
0	1	2	3
$\Phi(x, p_T^2)$	$\pi C_G^{[U]} \Phi_G$	$\pi^2 C_{GG,c}^{[U]} \Phi_{GG,c}$	$\pi^3 C_{GGG,c}^{[U]} \Phi_{GGG,c}$
$\tilde{\Phi}_\partial$	$\pi C_G^{[U]} \tilde{\Phi}_{\{\partial G\}}$	$\pi^2 C_{GG,c}^{[U]} \tilde{\Phi}_{\{\partial GG\},c}$	\dots
$\tilde{\Phi}_{\partial\partial}$	$\pi C_G^{[U]} \tilde{\Phi}_{\{\partial\partial G\}}$	\dots	\dots
$\tilde{\Phi}_{\partial\partial\partial}$	\dots	\dots	\dots

Table 2. TMD PDFs for an unpolarized or spin 0 target assigned to the quark correlators as given in Table 1.

PDFs FOR SPIN 0 HADRONS		
f_1	h_1^\perp	

Table 3. Assignments of TMD PDFs for a polarized spin 1/2 target.

PDFs FOR SPIN 1/2 HADRONS		
g_1, h_1	f_{1T}^\perp	$h_{1T}^{\perp(B1)}, h_{1T}^{\perp(B2)}$
g_{1T}, h_{1L}^\perp		
$h_{1T}^{\perp(A)}$		

Table 4. Assignments of TMD PFFs for unpolarized or spin 0 hadrons.

PFFs FOR SPIN 0 HADRONS		
D_1		
H_1^\perp		

Table 5. Assignments of TMD PFFs for polarized spin 1/2 hadrons.

PFFs FOR SPIN 1/2 HADRONS		
G_1, H_1		
$G_{1T}, H_{1L}^\perp, D_{1T}^\perp$		
H_{1T}^\perp		

5. Conclusions

We have introduced quark TMD correlators of definite rank in Eq. (22). In this new decomposition, we have made an expansion of the quark correlator into irreducible tensors multiplying correlators containing operator combinations of gluons, covariant derivatives and A -fields, the latter in the combination $i\partial = iD - gA$. In the decomposition gluonic pole factors contain the gauge link dependence, which are calculated from the transverse moments. These factors also give the process dependence, which is determined by the gauge link structure. The correlators of definite rank, in turn are parameterized in terms of the universal TMD PDFs depending on x and p_T^2 , such as given by Eqs. (23)-(29). The process dependence for a particular TMD PDF is in the same gluonic pole factors that appear in the expansion in Eq. (22).

An analysis for a quark spin 1/2 target shows that the process dependence is not strictly confined to the T-odd functions, such as the Sivers or the Boer-Mulders functions. In fact, we have shown the existence of three T-even pretzelosity functions. For fragmentation the TMD PFFs are already universal since gluonic pole matrix elements vanish for fragmentation correlators. Future work will be focused

on the study of universality for higher spin targets and gluon TMD PDFs^{22,25}. While for a spin 1/2 target one has at most rank two TMDs, one has for higher spins and gluon TMDs also higher rank functions, while also the color and gauge link structure is richer.

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